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# THE $OSp(1|4)$ SUPERPARTICLE AND EXOTIC BPS STATES<sup>a</sup>

Igor Bandos

*Institute for Theoretical Physics, National Science Center “Kharkov Institute  
of Physics and Technology”, Kharkov, 61108, Ukraine*

Jerzy Lukierski<sup>b</sup>

*Institute for Theoretical Physics University of Wrocław,  
50-204 Wrocław, Poland*

and

Dmitri Sorokin <sup>c</sup>

*INFN, Sezione Di Padova  
Via F. Marzolo, 8, 35131 Padova, Italia*

We discuss the dynamics of a superparticle in a superspace whose isometry is generated by the superalgebra  $OSp(1|4)$  or its central-charge contraction. Extra coordinates of the superspace associated with tensorial central charges are shown to describe spin degrees of freedom of the superparticle, so quantum states form an infinite tower of (half)-integer helicities. A peculiar feature of the model is that it admits BPS states which preserve  $3/4$  of target-space supersymmetries.

In this contribution we present results of work done on studying superparticle models whose symmetry and physical properties are defined by a supersymmetric algebra extended by tensorial charges <sup>1,2,3</sup>.

A motivation for this study has been to understand the physical meaning of tensorial “central” charges when they are associated not with superbranes but with relativistic (point-like) particles.

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<sup>c</sup>On leave from Kharkov Institute of Physics and Technology, Kharkov, 61108, Ukraine

As we shall see, in the case of superparticles tensorial central charges have different physical meaning than that of superbranes. They correspond to spin degrees of freedom of the superparticles, while it is well known that in the case of branes tensorial charges describe the coupling of branes to tensor gauge fields of target-space supergravity. That is, brane tensorial charges are similar to electric and magnetic charges of particles.

For instance, a membrane in  $D = 4$  or  $D = 11$  target space couples to a three-form gauge field  $A_{mnp}(x)$ . In the membrane worldvolume action <sup>4</sup> this coupling is described by the Wess-Zumino term

$$S_{coupl} = \int_{\mathcal{M}_3} dx^m dx^n dx^p A_{mnp}(x), \quad (1)$$

where the integral is taken over the membrane worldvolume  $\mathcal{M}_3$  of the pull back of the gauge field. As was shown in <sup>5</sup>, the membrane charge associated with this coupling is a two-rank tensor of the form

$$Z^{mn} = \int_{\mathcal{M}_2} dx^m \wedge dx^n, \quad (2)$$

where the integral is taken over the two-dimensional surface of the membrane.

Using the supermembrane action <sup>4</sup> the authors of <sup>5</sup> derived the form of the superalgebra generated by the Noether charges of the supermembrane. The supertranslation algebra thus obtained was shown to contain the contribution of the membrane charge (2) to the r.h.s. of the anticommutator of the fermionic supercharges

$$\{Q_\alpha, Q_\beta\} = -2(C\gamma_m)_{\alpha\beta} P^m + (C\gamma_{mn})_{\alpha\beta} Z^{mn}, \quad (3)$$

where  $P_m$  is the standard momentum (or bosonic translation) generator.

Thus, the worldvolume actions for the superbranes imply that the underlying supersymmetry of superbranes is described by a superalgebra extended by tensorial central charges <sup>5,6,7</sup>.

Recent analysis carried out in <sup>8,9</sup> has demonstrated that when superbranes propagate in anti-de-Sitter superbackgrounds, the superalgebra of their Noether supercharges gets extended to a corresponding maximal  $OSp$  superalgebra.

For instance, for branes in  $D = 10$  and  $D = 11$  target superspaces with  $AdS$  geometry the Noether charges generate the  $OSp(1|32)$  superalgebra. In <sup>10</sup> it has been assumed that the underlying superalgebra of M-theory should be even larger, namely,  $OSp(1|64)$  which is the minimal simple superalgebra which contains the supertranslation algebra of M-theory with a two-form and a five-form central charge as a *subsuperalgebra*.

If the  $OSp(1|32)$  and  $OSp(1|64)$  supergroup are related to M-theory, it seems instructive to find and study simple dynamical systems whose properties would be governed by  $OSp(1|2n)$  supergroups. And this has been another motivation for our work.

To simplify consideration let us take the  $OSp(1|4)$  supergroup which is the isometry of  $N = 1, D = 4$  superspace having the four-dimensional  $AdS_4$  space as a bosonic subspace (see <sup>3</sup> for the generic  $OSp(1|2n)$  case). The  $OSp(1|4)$  superalgebra has the following form

$$[M_{ab}, M_{cd}] = i(\eta_{ad}M_{bc} + \eta_{bc}M_{ad} - \eta_{ac}M_{bd} - \eta_{bd}M_{ac}), \quad (4)$$

$$[M_{ab}, P_c] = i(\eta_{bc}P_a - \eta_{ac}P_b), \quad (5)$$

$$[P_a, P_b] = \frac{i}{R^2}M_{ab}, \quad (6)$$

$$\begin{aligned} \{Q_\alpha, Q_\beta\} &= -2(C\gamma^a)_{\alpha\beta}P_a + \frac{1}{R}(C\gamma^{ab})_{\alpha\beta}M_{ab}, \\ [M_{ab}, Q_\alpha] &= -\frac{i}{2}Q_\beta(\gamma_{ab})^\beta_\alpha, \quad \gamma_{ab} = \frac{1}{2}(\gamma_a\gamma_b - \gamma_b\gamma_a), \end{aligned} \quad (7)$$

$$[P_a, Q_\alpha] = -\frac{i}{2R}Q_\beta(\gamma_a)^\beta_\alpha. \quad (8)$$

The generators  $M_{ab}$  form the  $SO(1, 3)$  subalgebra (4), and  $M_{ab}$  and  $P_a$  ( $a=0,1,2,3$ ) generate the  $SO(2, 3) \sim Sp(4)$  subalgebra of  $OSp(1|4)$ .  $Q_\alpha$  are four Majorana spinor generators of  $OSp(1|4)$ . The parameter  $R$  is the  $AdS_4$  radius, and  $C_{\alpha\beta}$  is a charge conjugation  $Sp(4)$ -invariant matrix.

The  $N = 1, D = 4$  supertranslation algebra with tensorial central charges (3) is a contraction of the  $OSp(1|4)$  superalgebra carried out in the following way. One rescales the  $SO(1, 3)$  generators as

$$M_{ab} = RZ_{ab} \quad (9)$$

and takes the limit  $R \rightarrow \infty$ , which corresponds to the limit where  $AdS_4$  becomes asymptotically flat. Then  $Z_{ab}$  and  $P_a$  become Abelian and commute with  $Q_\alpha$ , and the anticommutator of  $Q_\alpha$  reduces to that in (3).

Consider now a simple ‘‘oscillator’’ realization of  $OSp(1|4)$ . For this let us combine  $P_a$  and  $M_{ab}$  into a symmetric spin-tensor generator

$$M_{\alpha\beta} = -2(C\gamma_a)_{\alpha\beta}P^a + (C\gamma_{ab})_{\alpha\beta}Z^{ab} \quad (10)$$

so that

$$\begin{aligned} [M_{\alpha\beta}, M_{\gamma\delta}] &= -\frac{4i}{R}[C_{\gamma(\alpha}M_{\beta)\delta} + C_{\delta(\alpha}M_{\beta)\gamma}], \\ [M_{\alpha\beta}, Q_\gamma] &= -\frac{4i}{R}C_{\gamma(\alpha}Q_{\beta)}, \quad \{Q_\alpha, Q_\beta\} = M_{\alpha\beta}. \end{aligned} \quad (11)$$

If we introduce a Grassmann-even spinor operator  $\lambda_\alpha$ , which forms the Heisenberg algebra

$$[\lambda_\alpha, \lambda_\beta] = \frac{2i}{R}C_{\alpha\beta}, \quad (12)$$

and a Grassmann-odd scalar  $\psi$  such that

$$\psi^2 = \frac{1}{2}, \quad \psi\lambda_\alpha - \lambda_\alpha\psi = 0, \quad (13)$$

we can realize  $M_{\alpha\beta}$  and  $Q_\alpha$  as follows

$$M_{\alpha\beta} = \frac{1}{2}(\lambda_\alpha\lambda_\beta + \lambda_\beta\lambda_\alpha), \quad Q_\alpha = \lambda_\alpha\psi. \quad (14)$$

Using the commutation relations for  $\lambda_\alpha$  and  $\psi$  it is not hard to check that  $Q_\alpha$  and  $M_{\alpha\beta}$  represented in this way generate the  $OSp(1|4)$  superalgebra,  $\mathcal{Z}_A \equiv (\lambda_\alpha, \psi)$  forming the fundamental representation of  $OSp(1|4)$ .

Note that in the limit  $R \rightarrow \infty$ ,  $\lambda_\alpha$  become commuting quantities (see (12)). Then we observe that

$$P_a = \lambda\gamma_a\lambda \quad (15)$$

in which one recognizes the Cartan–Penrose relation implying that

$$P_a P^a \equiv 0, \quad (16)$$

due to  $D = 4$   $\gamma$ -matrix identities. We thus assume that at  $R \rightarrow \infty$  the “oscillator” realization of  $OSp(1|4)$  may correspond to a massless  $D = 4$  superparticle with  $\lambda_\alpha$  playing the role of a twistor-like variable.

Before contraction such a superparticle propagates in the supergroup manifold  $OSp(1|4)$  parametrized by coordinates  $x^a, y^{ab}$  and  $\theta^\alpha$  associated, respectively, with the generators  $P_a, M_{ab}$  and  $Q_\alpha$ .

An action for this superparticle can be constructed in a way similar to that used for the first time by Ferber<sup>11</sup> for developing the supertwistor formulation of supersymmetric field theories.

In our case the superparticle worldline  $Z^M(\tau) = (x^a(\tau), y^{ab}(\tau), \theta^\alpha(\tau))$  on the  $OSp(1|4)$  manifold is parametrized by the time variable  $\tau$ . To construct

the action we pick the worldline pull back of left-invariant  $OSp(1|4)$  Cartan one-forms taking values in the  $SO(2, 3)$  subalgebra of  $OSp(1|4)$

$$\Omega^{\alpha\beta}(Z) = dZ^M(\tau)\Omega_M^{\alpha\beta}(Z) \quad (17)$$

and contract the spinor indices with *commuting* spinor variables  $\lambda_\alpha(\tau)$  which are classical counterparts of the spinor operators (12) used to realize the generators (14) of  $OSp(1|4)$ .  $\lambda_\alpha$  become non-commutative upon solving for second-class constraints of the model and passing from Poisson to Dirac brackets.

The action we thus obtain has the twistor-like form

$$S = \int d\tau \lambda_\alpha \lambda_\beta \partial_\tau Z^M \Omega_M^{\alpha\beta}. \quad (18)$$

By construction (18) is invariant under the  $OSp(1|4)$  transformations of the coordinates  $Z^M = (x^a, y^{ab}, \theta^\alpha)$  and under the  $SO(2, 3)$  subgroup of  $OSp(1|4)$  acting in the tangent space of the supermanifold  $OSp(1|4)$ . To better understand the symmetry structure of the action (18) we note that the isometry of the supergroup manifold  $OSp(1|4)$  is the direct product of two supergroups  $OSp(1|4)_L \times OSp(1|4)_R$ . By using in (18) the  $OSp(1|4)$ -left-invariant Cartan forms corresponding only to  $SO(2, 3)_R$  we break “right-acting”  $OSp(1|4)_R$  down to  $SO(2, 3)$ .

To be able to analyze the action (18) one should know an explicit expression for the Cartan forms  $\Omega^{\alpha\beta}$ . This can be obtained by substituting an appropriate parametrization of the  $OSp(1|4)$  supergroup element  $G(Z)$  into the definition of the Cartan forms

$$\Omega = -iG^{-1}dG \equiv \Omega^{\alpha\beta}(Z)M_{\alpha\beta} + E^\alpha(Z)Q_\alpha, \quad (19)$$

where  $E^\alpha$  is a spinorial Cartan form associated with the supercharge generators.

In <sup>3</sup> we have found a parametrization of  $G(Z)$  which allowed us to obtain simpler expressions for the Cartan forms of  $OSp(1|4)$ , and generically of  $OSp(1|2n)$ , than those derived in earlier papers <sup>12,13,14</sup>. We have got

$$\Omega^{\alpha\beta} = v_\gamma^\alpha(y)v_\delta^\beta(y) \left[ (\gamma_a)^{\gamma\delta} e^a(x) + R(\gamma_{ab})^{\gamma\delta} \omega^{ab}(x) + \theta^{(\gamma} \mathcal{D}\theta^{\delta)} + (dvv^{-1})^{\gamma\delta} \right], \quad (20)$$

where  $v_\gamma^\alpha(y)$  are  $SO(1, 3)$  matrices in the spinor representation,  $e^a(x)$  and  $\omega^{ab}(x)$  are, respectively, the vierbein and the spin connection on the bosonic coset space  $AdS_4 = \frac{SO(2,3)}{SO(1,3)}$ , and  $\mathcal{D} = d + \frac{1}{2}\omega^{ab}(x)\gamma_{ab} + \frac{1}{2R}e^a(x)\gamma_a$  is the covariant  $AdS_4$  differential.

Substituting (20) into the superparticle action (18) and making the redefinition of  $\lambda_\alpha(\tau)$

$$\Lambda_\alpha(\tau) = \lambda_\beta v^\beta_\alpha(y), \quad (21)$$

we rewrite the action in the following form

$$S = \int_{\mathcal{M}_\tau} \left[ \Lambda \gamma_a \Lambda (e^a - i\theta \gamma^a \mathcal{D}\theta) + \Lambda \gamma_{ab} \Lambda \left( R\omega^{ab} + \frac{i}{2} \theta \gamma^{ab} \mathcal{D}\theta + \text{tr}(dv v^{-1} \gamma^{ab}) \right) \right]. \quad (22)$$

Note that the  $SO(1,3)$  coordinates  $y^{ab}$  enter this action only through the last term  $dv v^{-1}$ . All other terms depend on the bosonic  $x^a$  and fermionic  $\theta^\alpha$  coordinates of the coset superspace  $\frac{OSp(1|4)}{SO(1,3)}$  whose bosonic subspace is  $AdS_4$ .

If in (22) we skip the terms with  $\Lambda \gamma_{ab} \Lambda$  we will get the action describing a massless superparticle propagating in the  $AdS_4$  superspace<sup>3</sup>.

On the other hand, if we take the limit  $R \rightarrow \infty$  then, as we have discussed above, the  $OSp(1|4)$  superalgebra gets contracted to the super-Poincare algebra with the tensorial central charge, and the  $OSp(1|4)$  superparticle action reduces to the action which describes a superparticle propagating in  $N = 1, D = 4$  flat superspace  $(x^a, \theta^\alpha)$  extended by tensorial coordinates  $y^{ab}$  associated now with central charge momentum generators  $Z_{ab}$  (9). The superparticle action takes the form

$$S = \int [\Lambda \gamma_a \Lambda (dx^a - i\theta \gamma^a d\theta) + \Lambda \gamma_{ab} \Lambda (dy^{ab} + i\theta \gamma^{ab} d\theta)]. \quad (23)$$

This action, which by construction obeys supersymmetry with tensorial charges, was proposed in<sup>1</sup>. The massless  $N = 1, D = 4$  superparticle described by this action possesses quite unusual features.

One of them is that the action is invariant under fermionic (so called  $\kappa$ -symmetry) transformations with *three* independent parameters. The  $\kappa$ -symmetry transformations have the following form

$$\begin{aligned} \delta_\kappa \theta^\alpha &= \kappa^I(\tau) \mu_I^\alpha \quad (I = 1, 2, 3) \\ \delta_\kappa x^a &= i\theta \gamma^a \delta_\kappa \theta, \quad \delta_\kappa y^{ab} = -i\theta \gamma^{ab} \delta_\kappa \theta, \end{aligned} \quad (24)$$

where  $\mu_I^\alpha$  are three linearly independent commuting spinors orthogonal to  $\lambda_\alpha$ , i.e.  $\lambda_\alpha \mu_I^\alpha = 0$ .

Remember, that standard superparticle and, in general, superbrane actions are invariant under  $\kappa$ -symmetry transformations with the number of independent parameters being half the number of the spinor coordinates of target superspace. So in  $N = 1, D = 4$  superspace the standard massless superparticle is invariant under *two* independent  $\kappa$ -symmetries. As was realized

in <sup>15</sup>, in the twistor-like formulation  $\kappa$ -symmetries can be made irreducible and traded for  $n = 2$  extended worldline supersymmetry with transformation properties of  $\theta^\alpha$  and  $x^a$  being

$$\delta\theta^\alpha = \epsilon_1(\tau)\lambda^\alpha + \epsilon_2(\tau)(\gamma_5\lambda)^\alpha, \quad \delta x^a = \theta\gamma^a\delta\theta. \quad (25)$$

One can easily observe the difference between the transformations (24) and (25). The latter explicitly contain  $\lambda^\alpha$ , while the former involve spinors orthogonal to  $\lambda^\alpha$ .

The invariance of the superbrane action with respect to local fermionic transformations implies that there exist superbrane configurations which preserve the number of target space supersymmetries which is equal to or less than the number of independent  $\kappa$ -symmetries. Such supersymmetric states saturate the Bogomol'nyi-Prasad-Sommerfeld energy bound. Thus in the case of the standard superbranes the number of unbroken target-space supersymmetries is not higher than  $\frac{1}{2}$  supersymmetry of the target-space vacuum.

For instance, in the case of the standard  $N = 1, D = 4$  massless superparticle only two of four target-space supersymmetries are unbroken. While in the case of the superparticle with tensorial central charges there are BPS superparticle configurations with *three* target-space supersymmetries, i.e.  $\frac{3}{4}$  of supersymmetry remain unbroken <sup>1</sup>.

Recently the possibility of the existence of exotic BPS brane configurations preserving more than  $\frac{1}{2}$  supersymmetry has been discussed in <sup>16</sup>. The superparticle model based on the action (23) is an example of such configurations.

This unusual property is also characteristic of the superparticle propagating on the whole supergroup manifold  $OSp(1|4)$  described by the action (22) <sup>3</sup>. The algebraic reason for such an exotic situation is that, as we have discussed above, the  $OSp(1|4)$  superalgebra and its central charge contraction is realized in such a way that the anticommuting Poisson brackets of the Noether supercharges  $Q_\alpha$  derived from (22) is equal to the  $\lambda_\alpha$ -spinor bilinears

$$\{Q_\alpha, Q_\beta\} = M_{\alpha\beta} = \lambda_\alpha\lambda_\beta. \quad (26)$$

The matrix  $\lambda_\alpha\lambda_\beta$  is degenerate and has the rank one. Hence only one of four supergenerators  $Q_\alpha$  (14) has nonzero anticommutator. To single out this supercharge, let us introduce a basis  $(\mu^\alpha, \mu_I^\alpha)$  ( $I = 1, 2, 3$ ) in the spinor space such that (compare with (24))

$$\mu^\alpha\lambda_\alpha = 1, \quad \mu_I^\alpha\lambda_\alpha = 0.$$

Then, in view of (13) and (14),

$$Q = \mu^\alpha Q_\alpha = \psi \quad \Rightarrow \quad Q^2 = \frac{1}{2}$$

corresponds to one broken supersymmetry and three supercharges

$$Q_I = \mu_I^\alpha Q_\alpha \quad \Rightarrow \quad \{Q_I, Q_J\} = \{Q_I, Q\} = 0$$

anticommute with themselves and with  $Q$ . Hence,  $Q_I$  act trivially on BPS superparticle states and correspond to three unbroken supersymmetries.

Another feature of the superparticle model with the tensorial charge coordinates is the physical meaning of these extra variables. As the Hamiltonian analysis and the quantization of this superparticle model have shown <sup>2</sup>, only one of the six tensorial charge coordinates  $y^{ab}$  is independent due to a large number of constraints. This coordinate takes discrete integer values  $n = 2s$  and labels half-integer and integer helicities ( $s = 0, \pm\frac{1}{2}, \pm 1, \dots, \infty$ ) of massless quantum states of the superparticle in  $D = 4$  space-time.

Let us consider this in more detail. In the Weyl representation of spinors

$$\lambda_\alpha = (\lambda_A, \bar{\lambda}^{\dot{A}}), \quad \theta_\alpha = (\theta_A, \bar{\theta}^{\dot{A}}), \quad A = 1, 2; \quad \dot{A} = 1, 2$$

the action (23) takes the form

$$S = \int \left[ \lambda_A \bar{\lambda}_{\dot{A}} (dx^{A\dot{A}} - i\theta^A d\bar{\theta}^{\dot{A}} + i d\theta^A \bar{\theta}^{\dot{A}}) + \lambda_A \lambda_B (dy^{AB} - i\theta^A d\theta^B) + c.c. \right], \quad (27)$$

where

$$x^{A\dot{A}} = x^a \sigma_a^{A\dot{A}}, \quad y^{AB} = y^{ab} (\sigma_{ab})^{AB}, \quad y^{\dot{A}\dot{B}} = y^{ab} (\bar{\sigma}_{ab})^{\dot{A}\dot{B}},$$

and  $\sigma_a^{A\dot{A}}$  are the Pauli matrices  $(\sigma_{ab})_A^B = 1/2i \left( \sigma_{aA\dot{B}} \tilde{\sigma}_b^{\dot{B}B} - (a \leftrightarrow b) \right)$ .

From (27) we obtain that the canonical momenta associated with  $D = 4$  coordinates  $x^a$  and tensorial charge coordinates  $y^{ab}$  are, respectively,

$$p_{A\dot{A}} = \lambda_A \bar{\lambda}_{\dot{A}} \quad (28)$$

$$Z_{AB} = \lambda_A \lambda_B, \quad Z_{\dot{A}\dot{B}} = \bar{\lambda}_{\dot{A}} \bar{\lambda}_{\dot{B}}.$$

Notice, in particular, that  $p^{A\dot{A}} p_{A\dot{A}} \equiv 0$ , and, hence, the superparticle is massless.

The Cartan–Penrose relation (28) establishes the correspondence between three independent components of lightlike  $p_a = \sigma_a^{A\dot{A}} p_{A\dot{A}}$  and three components of  $\lambda_A, \bar{\lambda}_{\dot{A}}$ . Only the phase of  $\lambda$

$$\lambda_A = e^{i\varphi(\tau)} \lambda_A^0 \quad (29)$$

remains undetermined.



If  $y^{ab} = 0$ , we deal with a twistor superparticle considered in <sup>11,17</sup>, its action being

$$S = \int \lambda_A \bar{\lambda}_{\dot{A}} (dx^{A\dot{A}} - i\theta^A d\bar{\theta}^{\dot{A}} + id\theta^A \bar{\theta}^{\dot{A}}). \quad (30)$$

In addition to all symmetries of the action (27), the action (30) is invariant under local  $U(1)$  transformations

$$\lambda_A \rightarrow e^{i\varphi(\tau)} \lambda_A, \quad \bar{\lambda}_{\dot{A}} \rightarrow e^{-i\varphi(\tau)} \bar{\lambda}_{\dot{A}}. \quad (31)$$

This gauges away the phase component of  $\lambda_A$  and establishes the one-to-one correspondence between the independent components of the twistor superparticle momentum  $p_a$  and  $\lambda_A$ .

As the quantization of the action (30) has shown <sup>11,17</sup>, the quantum states of the  $N = 1, D = 4$  twistor superparticle form chiral supermultiplets of physical states with helicity 0 and  $\frac{1}{2}$ . These supermultiplets are described by chiral superfields  $\Phi(x^a - i\theta\sigma^a\bar{\theta}, \theta_A)$ .

In the case of the superparticle model (27) with additional tensorial coordinates  $y^{ab}$  there is no local  $U(1)$  symmetry (31). The compact phase component of  $\lambda_A$  becomes a *physical* momentum degree of freedom which corresponds to a single independent component of tensorial charge momenta  $Z_{ab}$ .

So the superparticle wave function in the momentum representation now becomes

$$\Phi(p_a, \varphi, \theta^\alpha) = \sum_{n=0}^{\infty} [e^{in\varphi} \Phi_n(p_a, \theta^\alpha) + e^{-in\varphi} \bar{\Phi}_n(p_a, \theta^\alpha)]. \quad (32)$$

Integer  $n$  is associated with an independent discrete (or quantized) component of the central charge coordinates  $y^{ab}$ , which is the Fourier image of the compact phase momentum component  $\varphi$ . This resembles a “dual” Kaluza–Klein effect when instead of a spatial direction, compactified is the corresponding momentum coordinate of the phase space.

In the Lorentz-covariant form the first-quantized wave function of the superparticle in the momentum representation looks as follows <sup>2</sup>

$$\Phi = \sum_{n=0}^{\infty} \Phi^{\alpha_1 \dots \alpha_n}(p_m) \lambda_{\alpha_1} \dots \lambda_{\alpha_n} + \chi \sum_{n=0}^{\infty} \Psi^{\alpha_1 \dots \alpha_n}(p_m) \lambda_{\alpha_1} \dots \lambda_{\alpha_n}, \quad (33)$$

where  $\chi = \theta^\alpha \lambda_\alpha, \chi^2 \equiv 0$ . Each component of this series forms an irreducible representation of the Lorentz group  $SO(1, 3)$  and describes a massless state with integer or half-integer superhelicity depending whether  $n$  is even or odd.

Thus, the superparticle model with tensorial charges produces, upon quantization, an infinite tower of massless fields of higher spin. The structure of the wave function which describes higher-spin states is similar to that used in the

formalism developed by Vasiliev (for a recent review see <sup>18</sup>) to construct the theory of higher-spin fields. Hence, the model which we have briefly described can be assumed to be a classical counterpart of the field theory of higher spins.

To conclude, we have demonstrated that tensorial charges appearing in extensions of supertranslation algebras may have different meaning than that one got accustomed to in superbrane models. They may describe spinning degrees of freedom of a dynamical system. Superparticle models on supergroup manifolds  $OSp(1|2n)$  and their contractions to flat superspaces with tensorial coordinates may give rise to exotic BPS configurations which preserve more than  $\frac{1}{2}$  supersymmetry.

We have seen that the  $AdS_4$  space is an intrinsic part of the construction of the superparticle action (22) on  $OSp(1|4)$ , which upon the contraction describes free higher-spin states. It is well known that to switch on interactions of higher-spin fields one needs the space-time to be of  $AdS$  geometry <sup>18</sup>. So an interesting problem to study is whether the superparticle model on  $OSp(1|4)$  may help to make a progress in constructing the theory of interacting fields of higher spin.

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